

Testing lepton non-unitarity with the next generation of (Germanium-based) CE ν NS reactor experiments

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Coherent elastic neutrino-nucleus scattering (CE ν NS) has been experimentally confirmed using neutrinos from pion decay at rest, solar neutrinos and reactor antineutrinos. Future CE ν NS experiments will foreseeable lead to precision measurements which will be a powerful tool to search for new physics beyond the Standard Model. In this work, we investigate possible deviations from unitarity in the 3×3 leptonic mixing matrix that controls the propagation of active neutrinos. Such deviations may originate from the mixing with additional gauge singlet fermions and depending on their mass scale and mixing, the resulting phenomenology can differ substantially. We explore two well-motivated regimes: the *seesaw limit*, where the new fermions are heavy and kinematically inaccessible, leading to effective deviations from unitarity in the active sector; and the *light sterile limit*, where they are light enough to be produced and participate in neutrino propagation and scattering processes. We show how these scenarios modify both CE ν NS and elastic neutrino-electron scattering (E ν eS), and we present the corresponding sensitivity projections for a future CE ν NS reactor experiment obtained by upscaling the CONUS+ experiment, which reported the first observation of reactor CE ν NS. We identify the leading experimental systematics relevant for such an upscaling and demonstrate the resulting capability to probe TeV-scale new physics. Our results highlight the strong potential of CE ν NS to test the structure of the lepton sector and to search for physics beyond the Standard Model.

I. INTRODUCTION

The discovery of neutrino oscillations [1, 2] implies physics Beyond Standard Model (BSM), with massive and non-degenerate neutrinos being the simplest explanation. This has led to many projects and ideas mostly driven by experiments involving charged current (CC) processes aiming at improved determinations of neutrino masses and mixings. New physics may, however, manifest itself also in modified CC interactions and especially also in neutral current (NC) processes. It may also show up as unitarity violation as a consequence of new, gauge singlet fermions and the associated phenomena have been extensively explored [3–28]. Another potentially powerful tool are coherently enhanced NC processes, in particular *coherent elastic neutrino-nucleus scattering* - (CE ν NS), pioneered by Freedman in the 70s [29]. CE ν NS has been recently observed with pion decay-at-rest neutrinos [30–33], solar neutrinos [34] and reactor antineutrinos [35] which provides an interesting and complementary way to study neutrinos [36]. In this paper we analyze the potential of precision CE ν NS measurements achievable with an upscaling of the successful CONUS+ technology.

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Gauge singlet fermions are in general well motivated and three right-handed neutrinos are in fact one of the best ways to explain neutrino masses. This allows Dirac and Majorana mass terms leading to the famous seesaw mechanism. Diagonalization of the total mass matrix leads then usually to three light Majorana neutrinos, where their mixing matrix is unitary if the right-handed Majorana states are assumed to be ultra heavy. The right-handed states lead, however, to unitarity violations if they are not so heavy. Unitarity violations also show up if more gauge singlet fermions exist. The structure of both the CC and NC changes in all these cases in a non-trivial way [37, 38]. Indeed, the mixing matrix underlying the leptonic CC weak interaction that describes oscillations [39] deviates from unitarity, while the NC interaction of mass eigenstate neutrinos deviates from the unit matrix, with these two features inter-connected. Although the effect of non-unitary neutrino mixing was first discussed in the context of astrophysical neutrino propagation [40–42], it can be phenomenologically relevant in earth-bound experiments.

This happens in the context of genuine low-scale seesaw schemes, such as the inverse [43, 44] or the linear seesaw mechanism [45–47], leading to potentially sizable deviations from the conventional leptonic weak currents with unitary mixing. These corrections are expressed as power series in the parameter $\varepsilon = \mathcal{O}(Yv/M)$, where M is the mediator mass scale and v is the SM vacuum expectation value (vev). Although small, we stress that ε can be non-negligible within low-scale realizations of the seesaw. We call this scenario *the seesaw limit*, where the new states are heavy enough to decouple from low energy experiments while the mixing with active neutrinos remains sizable and not suppressed by neutrino mass. In what follows, we remain agnostic about the neutrino mass generation and consider also the *light sterile limit*, where the gauge singlets are light enough to be produced and potentially participate in the neutrino propagation and scattering.

Here we explore the potential of CE ν NS and elastic neutrino-electron scattering (E ν eS) to probe non-minimal charged-current and neutral-current weak interactions within the context of CE ν NS experiments using Germanium detector technology close to a nuclear power reactor. Motivated by the recent detection of CE ν NS by the CONUS+ experiment [35] and the established role Germanium technology occupies in this field [48–56] we want to access the future potential of such experiments. At the same time, further experiments using other detection technologies are catching up, among them: CONNIE [57], NEON [58], NUCLEUS [59], RED-100 [60], RELICS [61] and Ricochet [62]. More insightful data can also be expected from the beam side [63, 64] and dark matter direct detection experiments [34, 65], where solar neutrinos have already been observed. First data provided by these experiments already allowed for various CE ν NS investigations within and beyond the SM such as determinations of the Weinberg angle at these low energies, searches for non-standard neutrino interactions (NSI), light (mediator) particles or electromagnetic properties of the neutrino [66–74]. Our study complements recent theoretical works that investigate both SM and BSM scenarios using Germanium reactor data [75–79], and aims to shed light on the scale of neutrino mass generation.

The paper is organized as follows: In Section II we introduce the theoretical framework we apply throughout this work and show how CE ν NS and E ν eS are altered by non-unitarity effects. In Section III we analyze the potential of a future experiment which we choose to be an upscaling of the successful CONUS+ technology to larger detector masses. Our findings are presented in Section IV for the cases of light and heavy new physics and for selected experimental specifications. We summarize our findings and conclude with Section V.

II. THEORY PRELIMINARIES

We assume that the 3 active neutrinos mix with m gauge singlet fermion states. In this case, the most general charged current weak interaction of massive neutrinos is described by a rectangular matrix K [37], connecting 3 charged lepton mass states with $3 + m$ massive neutral states:

$$-\mathcal{L}_{CC} = \frac{g}{\sqrt{2}} W_\mu^+ \bar{\ell}_L \gamma^\mu K \nu + h.c. , \quad (1)$$

with $\bar{\ell}_L = (\bar{e}_L, \bar{\mu}_L, \bar{\tau}_L)$ and $\nu = (\nu_1, \nu_2, \dots, \nu_{3+m})^T$ the mass eigenstates of the charged leptons and neutrinos, respectively. K is a rectangular $3 \times (3 + m)$ matrix where - in the basis where the charged lepton mass matrix is diagonal, the upper blocks are simply the first 3 rows of the neutrino mixing matrix [38]. We can also define the relevant sub-block as

$$K = \begin{pmatrix} N & S \end{pmatrix} , \quad (2)$$

with N a (3×3) matrix, while S is $(3 \times m)$. Further $KK^\dagger = I_{3 \times 3}$ and thus $SS^\dagger = I - NN^\dagger$, but note that in general $P = K^\dagger K \neq I$. The matrix P , which is square $(3 + m) \times (3 + m)$, non-diagonal, non-unitary and projective ($P^2 = P$) parametrizes the neutral current interaction given by

$$-\mathcal{L}_{NC} = \frac{g}{2 \cos \theta_W} \bar{\nu} \gamma^\mu P_L K^\dagger K \nu = \frac{g}{2 \cos \theta_W} \bar{\nu} \gamma^\mu P_L P \nu . \quad (3)$$

This formalism applies to both heavy and light singlet states. In what follows, we will now distinguish between two relevant limits, depending on the mass of the new neutral leptons.

A. Heavy neutral leptons: the seesaw limit

We first assume $m_4, m_5, \dots \sim M \gg \Lambda_{EW}$, which we call the seesaw limit. While our phenomenological analysis remains model-agnostic, it is useful to recall what is theoretically expected. In seesaw scenarios, the hierarchy between the heavy singlet states and the light active ones allows for an expansion in a small parameter ε , which quantifies the departure from unitarity in the active sector. In canonical high-scale seesaws, such as the type-I mechanism, this parameter scales as $\varepsilon \sim \mathcal{O}(\sqrt{m_\nu/M}) \ll 1$ and therefore leads to unobservable effects. By contrast, in genuine low-scale realizations, paradigmatically the inverse or linear seesaws, the expansion parameter is $\varepsilon \sim \mathcal{O}(m_D/M)$, with $m_D = Yv$ the Dirac mass generated by the Yukawa interaction of coupling Y and v the Higgs vacuum expectation value. Assuming $Y = \mathcal{O}(1)$, any experimental bound on ε can thus be interpreted as a direct constraint on the heavy-mediator mass scale M . Current global fits to oscillation data [80] typically require $\varepsilon^2 \lesssim \mathcal{O}(10^{-2})$, corresponding to $M \gtrsim 1.7 \text{ TeV}$ for $Y \sim 1$.

If the energy of a given process is well below the mass of the heavy mediators, they cannot be produced and therefore do not participate in oscillation experiments. Then, effectively, only the first 3×3 blocks of K and P will play a role in the weak interactions, i.e. N in the charged current and $NN^\dagger$ in the neutral current. We can relate the order of these blocks with the seesaw expansion parameter $\varepsilon \sim \mathcal{O}(m_D/M)$ as

$$NN^\dagger \sim 1 - \mathcal{O}(\varepsilon^2) , \quad SS^\dagger \sim \mathcal{O}(\varepsilon^2) . \quad (4)$$

The resulting non-unitary matrix N can be parametrized following [3] as ¹

$$N = \begin{pmatrix} \alpha_{11} & 0 & 0 \\ \alpha_{21} & \alpha_{22} & 0 \\ \alpha_{31} & \alpha_{31} & \alpha_{33} \end{pmatrix} \cdot U. \quad (5)$$

Besides the 3×3 unitary matrix U used to describe neutrino mixing in the conventional unitary case, one has the triangular prefactor characterized by 3 real diagonal α_{ii} ($i = 1, 2, 3$), and 3 non-diagonal α_{ij} ($i \neq j$) which are complex. Note that the α_{ij} parameters have a direct interpretation in terms of the mixing angles between active neutrinos and the heavy singlet states. As an illustrative example, in the $3 + 1$ scheme with one heavy neutral lepton we find $\alpha_{ii} = \cos \theta_{i4}$ for the diagonal entries, while the off-diagonal ones are given by $\alpha_{ij} = \sin \theta_{i4} \sin \theta_{j4} e^{i(\phi_{i4} - \phi_{j4})}$, where θ_{i4} and ϕ_{i4} are the mixing angle and phases between the active and sterile sectors. Since the mixing angles are expected to be small, the diagonal terms are close to 1 and real while the off-diagonal entries are small and complex. The general expressions for an arbitrary number of heavy neutral leptons can be found in the appendix of [3]. This is a convenient and complete description of non-unitarity in the lepton sector. By construction, N and S must satisfy the relation $NN^\dagger = 1 - SS^\dagger$, hence $NN^\dagger \sim 1 - \mathcal{O}(\varepsilon^2)$. Explicitly,

$$NN^\dagger = \begin{pmatrix} \alpha_{11}^2 & \alpha_{11}\alpha_{21}^* & \alpha_{11}\alpha_{31}^* \\ \alpha_{11}\alpha_{21} & \alpha_{22}^2 + |\alpha_{21}|^2 & \alpha_{22}\alpha_{32}^* + \alpha_{21}\alpha_{31}^* \\ \alpha_{11}\alpha_{31} & \alpha_{22}\alpha_{32} + \alpha_{21}^*\alpha_{31} & \alpha_{33}^2 + |\alpha_{31}|^2 + |\alpha_{32}|^2 \end{pmatrix}, \quad (6)$$

from which one can read off the strength of the α_{ij} in terms of the small seesaw expansion parameter ε :

$$\alpha_{ii}^2 \sim 1 - \mathcal{O}(\varepsilon^2), \quad (7)$$

$$|\alpha_{ij}|^2 \sim \mathcal{O}(\varepsilon^4), \quad i \neq j. \quad (8)$$

One sees that the strength of the off-diagonal α parameters is suppressed relative to the deviations of the flavour-diagonal ones from their SM values. In zero-distance experiments, where neutrinos cannot oscillate from the source to the detector, the 0th order in the seesaw expansion corresponds to the unitary limit, the 1st order gives only diagonal flavour-conserving effects, while the genuine flavour-violating effects of non-unitarity only appear at 2nd order. Notice also that this behavior is consistent with the validity of the well-known triangle inequality $|\alpha_{ij}| \leq \sqrt{(1 - \alpha_{ii}^2)(1 - \alpha_{jj}^2)}$ [3].

From Eq. 7 a bound on $1 - \alpha_{ii}^2$ can be interpreted, up to $\mathcal{O}(1)$ factors, as a bound on ε^2 . Using $\varepsilon \sim m_D/M$ and $m_D = Yv$, one obtains the illustrative relation

$$M \gtrsim v \frac{1}{\sqrt{(1 - \alpha_{ii}^2)^{\text{limit}}}}. \quad (9)$$

Here $(1 - \alpha_{ii}^2)^{\text{limit}}$ denotes the experimental upper bound on the non-unitarity parameter $(1 - \alpha_{ii}^2)$. Eq. 9 should not be interpreted as a strict experimental constraint on M , but rather as an intuitive translation between the α parameters and the characteristic mass scale of low-scale seesaw mediators, assuming $\mathcal{O}(1)$ Yukawa couplings.

Additionally, unitarity violation leads to a redefinition of the Fermi constant, which is extracted from the muon lifetime assuming the SM to be valid. In the presence of non-unitarity the measured quantity would be

¹ An alternative description and its relationship with Eq. 5 is discussed in Ref. [81].

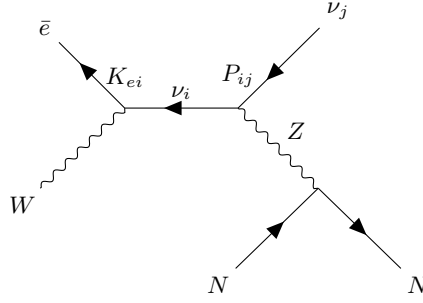


FIG. 1: Feynman diagram of CE ν NS. Modifications due to lepton non-unitarity introduce corrections at the neutrino interaction vertices.

an effective muon decay coupling G_μ . Since the W boson vertices are modified by the non-unitarity parameters one finds

$$G_\mu = 1.1663787(6) \times 10^{-5} \text{ GeV}^{-2} \quad (\text{effective } \mu^- \text{ decay constant [82]}) \quad (10)$$

$$G_\mu^2 = G_F^2 (NN^\dagger)_{ee} (NN^\dagger)_{\mu\mu} \quad (11)$$

and therefore

$$1 \leq \frac{G_F^2}{G_\mu^2} = \frac{1}{(NN^\dagger)_{ee} (NN^\dagger)_{\mu\mu}} \approx 3 - \alpha_{11}^2 - \alpha_{22}^2 \sim 1 + \mathcal{O}(\varepsilon^2) . \quad (12)$$

Consequently, any process proportional to G_F^2 will receive an “enhancement” due to the deviation from unitarity. This is counterintuitive because naively one expects less events than in the SM if the mixing is non-unitary, due to kinematically inaccessible heavy states. The reduction of the event number due to non-unitarity, and the “increase” due to the redefinition of G_F compete with each other, so that in some cases one can achieve $\mathcal{N}_{\text{SM}}/\mathcal{N}_{\text{NU}} = 1$ even in the presence of non-unitarity.

In the context of future CONUS-like experiments, we are interested in the electron antineutrinos produced by the reactor and the two accessible processes. The first is CE ν NS, $\bar{\nu}_e + \text{Ge} \rightarrow \bar{\nu}_j + \text{Ge}$, where N is a nucleus and the outgoing neutrino is not measured, so we will sum over the light mass eigenstates $\bar{\nu}_j$. The other is E ν eS, $\bar{\nu}_e + e \rightarrow \bar{\nu}_j + e$. The relevant diagrams are given in figs. 1 and 2, which include also the neutrino production vertices. In the unitary (SM) limit, the differential CE ν NS cross section is given by

$$\frac{d\sigma}{dT_A}(T_A, E_\nu) = \frac{G_F^2}{4\pi} m_A Q_W^2 \left(1 - \frac{m_A T_A}{2E_\nu^2} \right) |F(T_A)|^2, \quad \text{with } Q_W = [(1 - 4\sin^2 \theta_W)Z - N] . \quad (13)$$

Here, m_A is the nucleus mass and Q_W is the weak nuclear charge, which is defined by the number of protons Z and neutrons N , respectively. For the nuclear form factor, the Helm parameterization is used [83]:

$$F(T_A) = \frac{3j_1(q(T_A)R_1)}{q(T_A)R_1} \exp \left[-\frac{(q(T_A)s)^2}{2} \right], \quad (14)$$

with the spherical Bessel function j_1 , the momentum transfer $q^2 = 2m_A T_A$, the nuclear skin thickness $s \simeq 1$ fm, $R_1 = \sqrt{R^2 - 5s^2}$ and $R \simeq 1.2A^{1/3}$ fm. Due to small momentum transfers at a reactor site, the impact

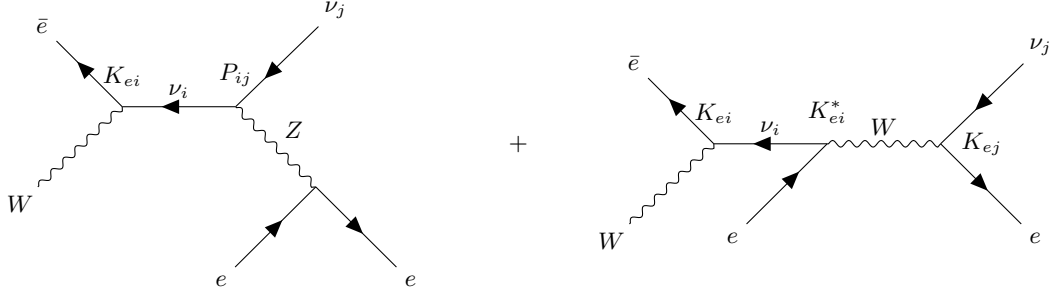


FIG. 2: Feynman diagrams for $E\nu eS$. In this case, there exist two diagrams with modifications depending on the mediator present in the interaction.

of the nuclear form factor is negligible, i.e. $F \rightarrow 1$. However, with increasing precision on the SM signal - as investigated in this work - this quantity becomes more relevant.

The target recoil energies T_x for $x = \{e, A\}$ depend on the scattering angle θ (lab frame) and is given by

$$T_x = \frac{2m_x E_\nu^2 \cos^2 \theta}{(m_x + E_\nu)^2 - E_\nu^2 \cos^2 \theta} \xrightarrow{\theta \rightarrow 0} \frac{2E_\nu^2}{m_x + 2E_\nu}, \quad (15)$$

where the last step defines the maximal nuclear recoil $T_x^{\max}$.

Compared to the SM, each vertex is modified in the presence of non-unitarity. As such, the probability factor attached to the diagram in 1 is given by

$$\mathcal{P} = N_{ei} N_{\alpha i}^* N_{\alpha j} N_{ek}^* N_{\beta k} N_{\beta j}^* = (NN^\dagger NN^\dagger NN^\dagger)_{ee}. \quad (16)$$

By taking into account the redefinition of the Fermi constant, performing the seesaw expansion and keeping terms up to order $O(\varepsilon^2)$ we find the expected number of $CE\nu NS$ events compared to the SM case to be

$$\left(\frac{\mathcal{N}_{\text{NU}}}{\mathcal{N}_{\text{SM}}} \right)^{\text{CE}\nu\text{NS}} = \mathcal{P} \frac{G_F^2}{G_\mu^2} = \frac{(NN^\dagger NN^\dagger NN^\dagger)_{ee}}{(NN^\dagger)_{ee} (NN^\dagger)_{\mu\mu}} \approx 2\alpha_{11}^2 - \alpha_{22}^2. \quad (17)$$

This ratio is equal to 1 in the unitary limit, where $\alpha_{11} = \alpha_{22} = 1$.

On the other hand, the differential cross section of the elastic neutrino scattering on electrons in the unitary limit is given by

$$\frac{d\sigma}{dT_e}(T_e, E_\nu) = \frac{G_F^2 m_e}{2\pi} \left[(g_V + g_A)^2 + (g_V - g_A)^2 \left(1 - \frac{T_e}{E_\nu} \right)^2 + (g_A^2 - g_V^2) \frac{m_e T_e}{E_\nu^2} \right], \quad (18)$$

with $g_V = -\frac{1}{2} + 2\sin^2 \theta_W$ and $g_A = -\frac{1}{2}$ for electron antineutrinos.

Now we need to compute the probability factors associated to the neutral and charged current as well as their interference. In general, they are different in the presence of non-unitarity. Therefore, the dependence of the differential cross section on the recoil energy of the final electron will change with respect to the SM. However,

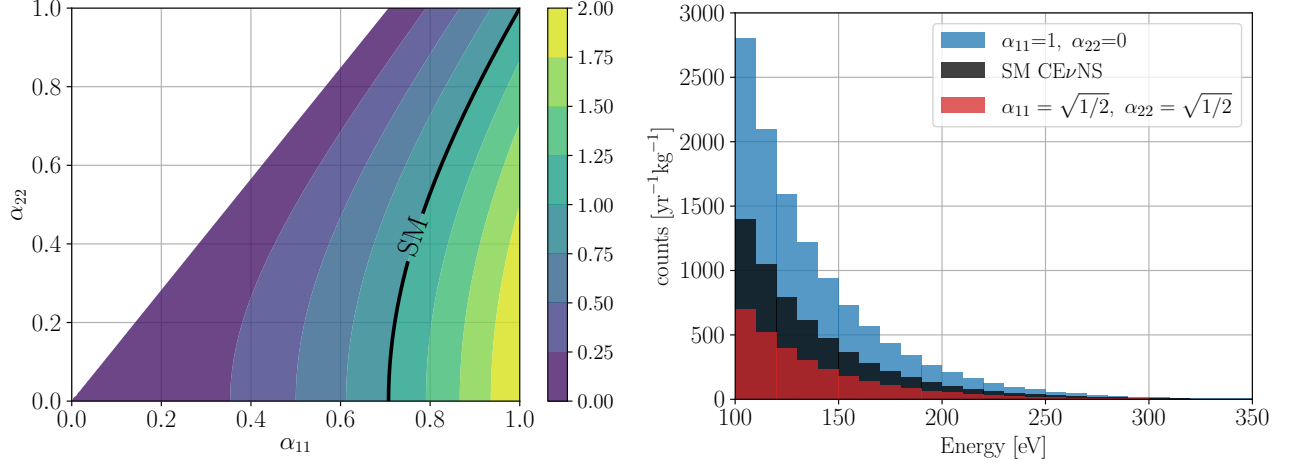


FIG. 3: Left: Prefactor $(2\alpha_{11}^2 - \alpha_{22}^2)$ present in Eqs. (17) and (23). Depending on the parameter configuration, the expected CE ν NS / E ν eS signal can be smaller or larger than the SM expectation indicated by the black line. Right: Exemplary CE ν NS spectra for the given alpha combinations, which are chosen to receive a bisection and doubling of the events. A flux of $\phi \sim 1.5 \cdot 10^{13}/\text{cm}^2/\text{s}$ is assumed ($L = 20 \text{ m}$, $P_{\text{th}} = 3.5 \text{ GW}$).

this deviation of the shape of the spectrum is not only be very hard to measure, but also 'theory suppressed' as we will see. Indeed, the probability factors are given by

$$\mathcal{P}_{NC} = N_{ei} N_{\alpha i}^* N_{\alpha j} N_{ek}^* N_{\beta k} N_{\beta j}^* = (NN^\dagger NN^\dagger)_{ee}, \quad (19)$$

$$\mathcal{P}_{CC} = N_{ei} N_{ei}^* N_{ej} N_{ek}^* N_{ek} N_{ej}^* = (NN^\dagger)_{ee}^3, \quad (20)$$

$$\mathcal{P}_{int} = N_{ei} N_{\alpha i}^* N_{\alpha j} N_{ek}^* N_{ek} N_{ej}^* = (NN^\dagger)_{ee} (NN^\dagger NN^\dagger)_{ee}, \quad (21)$$

which at order $O(\varepsilon^2)$ in the seesaw expansion become

$$\mathcal{P}_{NC} \approx \mathcal{P}_{CC} \approx \mathcal{P}_{int} \approx \alpha_{11}^6. \quad (22)$$

As a consequence

$$\left(\frac{\mathcal{N}_{\text{NU}}}{\mathcal{N}_{\text{SM}}} \right)^{\text{E}\nu\text{eS}} = \alpha_{11}^6 \frac{G_F^2}{G_\mu^2} = \frac{\alpha_{11}^6}{(NN^\dagger)_{ee} (NN^\dagger)_{\mu\mu}} \approx 2\alpha_{11}^2 - \alpha_{22}^2 = \left(\frac{\mathcal{N}_{\text{NU}}}{\mathcal{N}_{\text{SM}}} \right)^{\text{CE}\nu\text{NS}}, \quad (23)$$

which is the same modification as in the CE ν NS case.

B. Light sterile neutrinos

If instead the gauge singlet states are light enough to be produced and propagate, the full matrix K will appear in the description of the neutrino propagation, instead of only the sub-block N . It is also easy to see

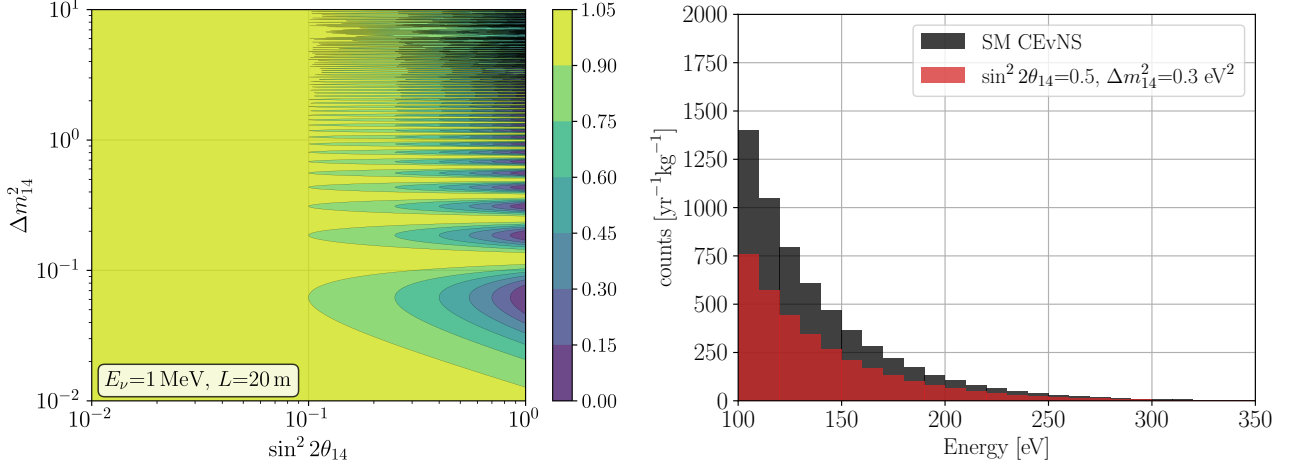


FIG. 4: Left: Oscillation probability of Eqs. (26) and (31) present in the CE ν NS and E ν eS cross section for fixed L/E_ν . Right: Expected CE ν NS spectrum of a light sterile neutrino for the given parameter combination compared to the SM case. Again, a flux of $\phi \sim 1.5 \cdot 10^{13}/\text{cm}^2/\text{s}$ is assumed ($L = 20$ m, $P_{\text{th}} = 3.5$ GW).

that, in this case, G_F does not get redefined, since the light sterile states leave the experiment undetected. For concreteness, we focus on the simple $3 + 1$ case, where we consider the addition of a 4th gauge singlet to the 3 massive active neutrinos. In doing so, we remain agnostic regarding the neutrino mass generation mechanism. In this scenario, the probability factor for the CE ν NS process including oscillation effects is given by

$$\mathcal{P} = K_{ei} e^{-iE_i t} K_{\alpha i}^* K_{\alpha j} K_{ek}^* e^{iE_k t} K_{\beta k} K_{\beta j}^* = (K e^{-iEt} K^\dagger K e^{iEt} K^\dagger)_{ee}, \quad (24)$$

where we used $KK^\dagger = 1$. In the basis where the charged lepton mass matrix is diagonal, K is simply the upper 3×4 rectangular block of the full 4×4 mixing matrix U . In general, U is parametrized by 6 angles and 6 CP-violating phases, making the above expression complicated and not enlightening. Since the mixing between the sterile neutrino and the active states is necessarily small, we can however expand in powers of the small mixing angles θ_{4i} and identify the rest with the standard mixing angles. Moreover, given the short baseline and the energy of reactor antineutrinos, we can neglect standard oscillation effects, i.e. $\Delta E_{ij} t \approx 0$ for $i, j \in \{1, 2, 3\}$, but we assume that $\Delta E_{4i} t$ is within experimental reach, i.e. oscillations from active to sterile states can occur. In this way, we recover the standard electron-electron oscillation survival probability in the presence of a light sterile neutrino at short distance, which only depends on one mixing angle and one mass squared difference:

$$\mathcal{P} \approx 1 - \sin^2 2\theta_{14} \sin^2 \left(\frac{L \Delta m_{41}^2}{4E_\nu} \right), \quad (25)$$

with the mixing angle $\sin^2 \theta_{14}$, experimental baseline L and the mass-squared difference Δm_{41}^2 . Thus, we obtain for CE ν NS

$$\left(\frac{\mathcal{N}_{\text{light}}}{\mathcal{N}_{\text{SM}}} \right)^{\text{CE}\nu\text{NS}} = 1 - \sin^2 2\theta_{14} \sin^2 \left(\frac{L \Delta m_{41}^2}{4E_\nu} \right). \quad (26)$$

Similarly, for $E\nu eS$ we compute the neutral, charged and interference prefactors being

$$\mathcal{P}_{NC} = (K e^{-iEt} K^\dagger K e^{iEt} K^\dagger)_{ee} \approx 1 - \sin^2 2\theta_{14} \sin^2 \left(\frac{L\Delta m_{41}^2}{4E_\nu} \right), \quad (27)$$

$$\mathcal{P}_{CC} = K_{ei} e^{-iE_i t} K_{ei}^* K_{ej} K_{ej}^* K_{ek}^* e^{iE_k t} K_{ek} = |(K e^{-iEt} K^\dagger)_{ee}|^2, \quad (28)$$

$$\mathcal{P}_{int} = K_{ei} e^{-iE_i t} K_{ki}^* K_{kj} K_{kj}^* K_{el}^* e^{iE_l t} K_{el} K_{ej}^* = (K e^{-iEt} K^\dagger K K^\dagger)_{ee} (K e^{iEt} K^\dagger)_{ee} = \mathcal{P}_{CC}. \quad (29)$$

We can simplify $\mathcal{P}_{CC}$ by again neglecting the standard oscillations and expanding in powers of the small quantities θ_{4i} and find

$$\mathcal{P} \equiv \mathcal{P}_{NC} \approx \mathcal{P}_{CC} \approx \mathcal{P}_{int} \approx 1 - \sin^2 2\theta_{14} \sin^2 \left(\frac{L\Delta m_{41}^2}{4E_\nu} \right), \quad (30)$$

which implies

$$\left(\frac{\mathcal{N}_{\text{light}}}{\mathcal{N}_{\text{SM}}} \right)^{E\nu eS} = 1 - \sin^2 2\theta_{14} \sin^2 \left(\frac{L\Delta m_{41}^2}{4E_\nu} \right) = \left(\frac{\mathcal{N}_{\text{light}}}{\mathcal{N}_{\text{SM}}} \right)^{CE\nu NS}. \quad (31)$$

III. EXPERIMENTAL FRAMEWORK AND ANALYSIS

Here we assume an upscaling of the Germanium-based detector technology used in CONUS+ that has recently achieved the first detection of $CE\nu NS$ with reactor antineutrinos [35]. This technology allows to go to larger detector masses which we consider up to 100 kg. In addition to a larger detector mass they have a low energy threshold [56] for which we also consider conceivable improvements. We assume as source nuclear power reactors that have a very intense electron-antineutrino flux. Specifically we assume the experiment to be located at a 20 m-distance to a typical commercial reactor with a thermal power of 3.5 GW. We choose a typical fuel composition of a pressurized water reactor (^{235}U , ^{238}U , ^{239}Pu , ^{241}Pu) = (56.1, 7.6, 30.7, 5.6) % and select a data-based reactor antineutrino spectrum. For this, we use the method proposed in [84] and utilize the provided unfolded spectra of inverse beta decay (IBD) measurements. The high energy part of the spectrum provided by [85] and calculated spectra for the energy region below the IBD threshold of [86] were added with appropriate normalization. Similar spectra have already been used in previous data analyses [35, 52]. This leads to an electron antineutrino flux of $\sim 1.5 \cdot 10^{13}/\text{cm}^2/\text{s}$ at the experimental site with underlying uncertainties - among others, of the reactor thermal power and fission fractions. A 3% level uncertainty can be considered established, but we will also show the impact of an improved uncertainty that could arise from combining all other existing and upcoming reactor experiments.

The cross sections of $CE\nu NS$ and $E\nu eS$ weighted with the spectral information are obtained via

$$\frac{d\sigma}{dT_x}(T_x) = \int_{E_{\min}}^{E_{\max}} dE_\nu \frac{dN}{dE_\nu}(E_\nu) \frac{d\sigma}{dT_x}(T_x, E_\nu), \quad (32)$$

with $T_{\{e^-, N\}}$ being the electron and nuclear recoil, respectively, E_ν the antineutrino energy and dN/dE_ν the reactor spectrum reaching up to $E_{\max} \simeq 11$ MeV. The minimal recoil energy is set by Eq. (15).

Further, we assume flat background levels of 10 cnts/keV/kg/d in the region of interest (ROI) for $CE\nu NS$ investigations, i.e. below 1 keV, and background contributions of 0.5 cnts/keV/kg/d above as reference values. More details about background events in reactor environments and individual contributions can be found in

[49, 87]. To reduce uncertainties on the background, we fit times of operating (ON) and shut down (OFF) reactor simultaneously and assume $t_{\text{OFF}} = 0.1 \cdot t_{\text{ON}}$.² Effects of overall improved background levels are discussed below. In this study we assume experimental exposures of (5, 50, 500) kg·yrs (reactor ON). 5 kg·yrs corresponds to CONUS+ operation which we term as “now”. 50 kg·yrs is something that can be obtained by operating CONUS+ with the latest upgraded detectors for a few years, which we call “soon”. 500 kg·yrs corresponds to an upscaling with 100 kg detector mass operated for 5 years, which we call “future”.

For the Germanium detectors we assume 100% detection efficiency in the ROI down to their threshold energies, for which we choose a (“now”, “soon”, “future”) value of (150, 125, 100) eV. Since this detector type is only sensitive to ionization, we need to account for the conversion of nuclear recoils T into charge signals E . The relevant quenching factor is so far well described by the widely applied Lindhard model [88] that has been confirmed to be valid for Germanium semiconductor detector down to the energies of interest. We use for the theory’s k parameter (reflecting the ratio between ionization and recoil energy at 1 keV) the measured value of $k = 0.162 \pm 0.004$ [89]. Future measurements may reduce the uncertainty on the k parameter and we will discuss the impact of such improvements. The conversion from recoil to ionization energy is done by a variable transform

$$\frac{d\sigma}{dE}(E) = \left[Qf^{-1}(E) + E \frac{dQf^{-1}}{dE}(E) \right] \frac{d\sigma}{dT_N}(Qf^{-1}(E) \cdot E), \quad (33)$$

with $Qf^{-1}(E)$ being the inverted Lindhard model in terms of ionization (detected) energy E .

Furthermore, we assume a connection between the detectors threshold energy and the intrinsic noise of the read-out electronic. In particular, we impose the detector threshold to be three times the FWHM (full width at half maximum) of an artificial test pulse at zero energy $E_{\text{thr}} \sim 3 \cdot \text{FWHM}$. As a consequence, the detector resolution is described by a Gaussian with an energy-dependent width given by

$$\sigma^2(E) = \left(\frac{E_{\text{thr}}}{3 \cdot 2.355} \right)^2 + \epsilon \cdot F \cdot E, \quad (34)$$

with the energy necessary to create an energy-whole pair in Germanium $\epsilon = 2.96$ eV (at 90 K) and the so-called Fano factor $F = 0.11$.

Taking all these aspects into account, the expected SM events rate are (16, 9, 5) cnts/kg/d for $\text{CE}\nu\text{NS}$ in the region (100, 125, 150) eV up to 1 keV and ~ 1.4 cnts/kg/d/keV for $\text{E}\nu\text{eS}$ up to 100 keV. Exemplary $\text{CE}\nu\text{NS}$ spectra for the scenarios under consideration can be found in the right plots of the figures 3 and 4, respectively.

Our sensitivity estimates are obtained with a Likelihood function that incorporates both reactor ON and reactor OFF data and pull terms taking into account the experimental uncertainties of the neutrino flux $\Delta\Phi$ and the Lindhard model Δk ,

$$-2\log \mathcal{L} = -2\log \mathcal{L}_{\text{ON}} - 2\log \mathcal{L}_{\text{OFF}} + \text{pull terms}. \quad (35)$$

Two model parameters - $(\alpha_{11}, \alpha_{22})$ in the seesaw and $(\sin^2 2\theta_{14}, \Delta m_{41}^2)$ in the light sterile limit - are fit together with two background normalization parameters $b_{<1 \text{ keV}}$ and $b_{>1 \text{ keV}}$, while a 3% and 1% uncertainties were assumed for the reactor antineutrino flux $\Delta\Phi$ and quenching given by the Lindhard model Δk , respectively. In addition, we allow the Weinberg angle $\sin^2 \theta_W$ to vary within current uncertainties at low energy $\sin^2 \theta_W =$

² Nuclear power reactors typically run eleven months followed by a month for re-fueling.

0.2374 ± 0.0020 [90]. Additionally, we determine results for a factor 10 improvement in $\Delta\Phi$ and the assumed background levels as well as a factor 2 in Δk to quantify the impact of these parameters. In doing so, we underline which experimental parameter are worth improving for future experiments. We perform a likelihood ratio test and extract limits (at 90% C.L.) on the parameter space from a χ^2 -distributed test statistics with two degrees of freedom. For the seesaw limit, we incorporate knowledge extracted from oscillations experiments [80] by adding an additional two-dimensional pull term to the likelihood function above. In doing so, we underline the potential of future CE ν NS reactor experiments in 'global fits' when data from several neutrino experiments are combined.

IV. RESULTS

A. Seesaw limit

As can anticipated from the count rates mentioned in the previous section, our limits are mainly driven by the CE ν NS channel. In addition, this has been confirmed by investigating the impact of both interaction channels individually, cf. figure 13. Although there exist already strong constraints from oscillation experiments [80], an improvement for α_{11} is expected due to an appearing factor of 2 in the prefactors of Eqs. (17) and (23).

At first we assess the sensitivity to the individual non-unitarity parameters by fixing the other one to unity. Our results are summarized in figure 5 and table I for our reference configuration and figure 6 and table II for optimized experimental features, respectively. Detailed $\Delta\chi^2$ profiles for the first, also when combined with oscillation data, can be found in the appendix, cf. figure 10. We note that limits improve with better detection thresholds and increased exposure. However, advances are stronger for the transition from recent to soon available thresholds and the first increase in exposure. For the highest assumed exposure, there is no clear improvement regardless of the chosen detector threshold, cf. figure 5. This indicates that the assumed systematic uncertainties, i.e. the antineutrino flux and signal quenching, become dominant. A general improvement is obtained when our reactor-only analysis is combined with information from oscillation experiments. While for α_{22} existing bounds from oscillation experiments are already quite strong, limits on α_{11} improve in a combined analysis, cf. table I, clearly underlining the importance of combined approaches in the future.

Looking at the configuration with reduced uncertainties and lower background level (a factor 10 for flux and background and a factor 2 for quenching), we obtain limits a factor > 2 better than for the previous case, cf. table II. Especially the larger exposures benefit from these optimization and yield increasingly better constraints. When combined with knowledge from oscillation, limits further improve with lower detector threshold and larger exposure. The obtained limits can be converted into an approximate mass scale where the connected new particles are expected to appear, cf. Eq. (9), assuming a low-scale seesaw and $\mathcal{O}(1)$ parameters. In doing so, we could soon (50 kg·yr exposure and 125 eV-threshold) constrain new physics to lie above ~ 1100 GeV and ~ 760 GeV for α_{11} and α_{22} , respectively. An improved setup could lift these "bounds" up to ~ 1900 GeV (α_{11}) and ~ 1400 GeV (α_{22}).³ An optimistic scenario with a 100 eV-threshold and 500 kg·yr combined with oscillations and the optimized experimental setup would constrain new physics up to 2500 GeV.

³ Here, we chose the intermediate values for exposure and detector threshold to show the experimental potential that is going to be available in the near future.

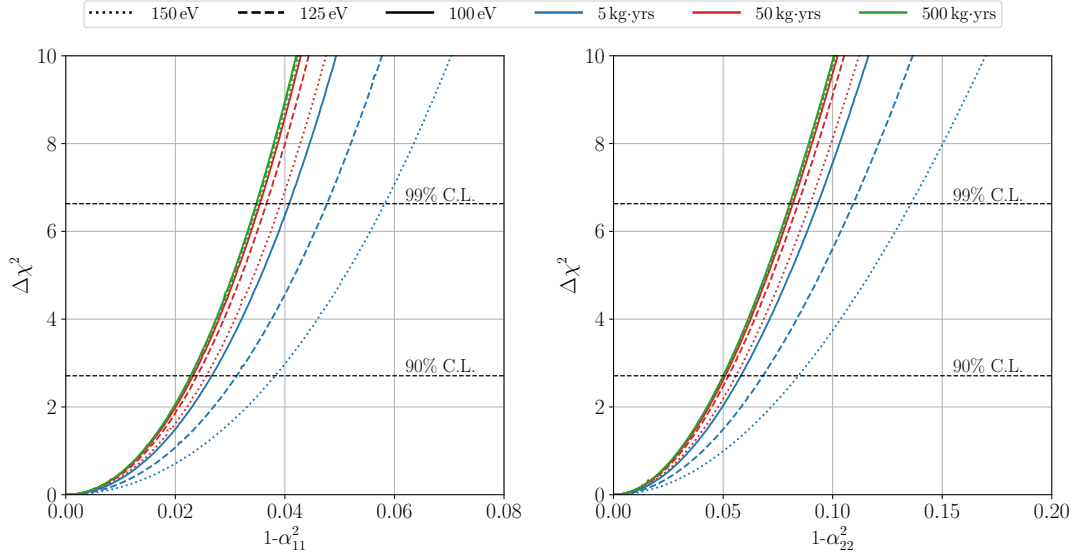


FIG. 5: $\Delta\chi^2$ profiles of the individual alpha parameters - the other one being fixed to unity - for the assumed experimental configuration. We show three threshold configurations together with three assumptions on the experimental exposure, which are indicated by different colors and line styles. Detailed profiles for the individual thresholds also combined with knowledge from oscillation experiments are illustrated in figure 10 in the appendix.

		$1 - \alpha_{11}^2$		$1 - \alpha_{22}^2$	
		+ osci.		+ osci.	
150 eV	5 kg.yr	0.039	0.022	0.086	0.008
	50 kg.yr	0.026	0.018	0.057	0.008
	500 kg.yr	0.023	0.016	0.051	0.008
125 eV	5 kg.yr	0.032	0.020	0.069	0.008
	50 kg.yr	0.024	0.017	0.053	0.008
	500 kg.yr	0.023	0.016	0.051	0.008
100 eV	5 kg.yr	0.027	0.018	0.059	0.008
	50 kg.yr	0.024	0.016	0.052	0.008
	500 kg.yr	0.023	0.016	0.051	0.008
oscillations (90% CL) [80]		0.061		0.01	

TABLE I: Individual 90% C.L. limits on the alpha parameters ($1 - \alpha_{ii}^2$) (the other one fixed to unity) for the assumed experimental configuration. The left columns show the limits for our experimental configuration alone, while right columns indicate limits obtained when knowledge from oscillation experiments is incorporated. The corresponding (oscillation) limits are given in the bottom row.

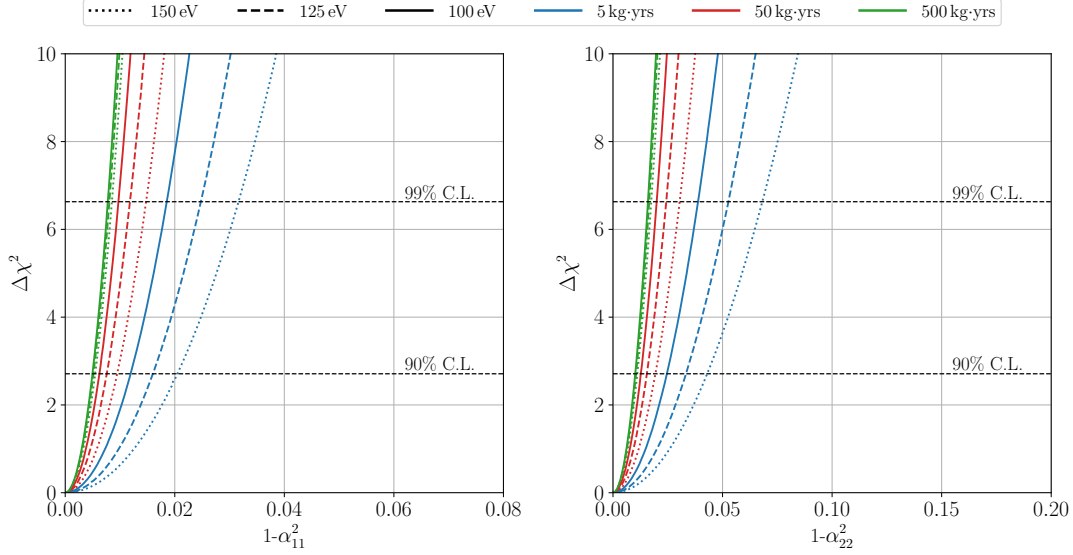


FIG. 6: $\Delta\chi^2$ profiles of the individual alpha parameters - the other one being fixed to unity - for the optimized experimental configuration. Again, we show three threshold configurations together with three experimental exposures, which are indicated by different colors and line styles.

		$1 - \alpha_{11}^2$		$1 - \alpha_{22}^2$	
		+ osci.		+ osci.	
150 eV	5 kg-yr	0.021	0.015	0.044	0.008
	50 kg-yr	0.010	0.008	0.020	0.007
	500 kg-yr	0.006	0.005	0.011	0.006
125 eV	5 kg-yr	0.016	0.013	0.034	0.008
	50 kg-yr	0.008	0.007	0.016	0.007
	500 kg-yr	0.005	0.005	0.011	0.006
100 eV	5 kg-yr	0.012	0.010	0.025	0.008
	50 kg-yr	0.006	0.006	0.013	0.007
	500 kg-yr	0.005	0.005	0.010	0.006
oscillations (90% CL) [80]		0.061		0.01	

TABLE II: Individual 90% C.L. limits on the alpha parameters ($1 - \alpha_{ii}^2$) (the other one fixed to unity) for the optimized experimental configuration. The left columns shows the limits from our experimental configuration alone, while right columns indicate limits obtained when knowledge from oscillation experiments is incorporated. The corresponding (oscillation) limits are given in the bottom row.

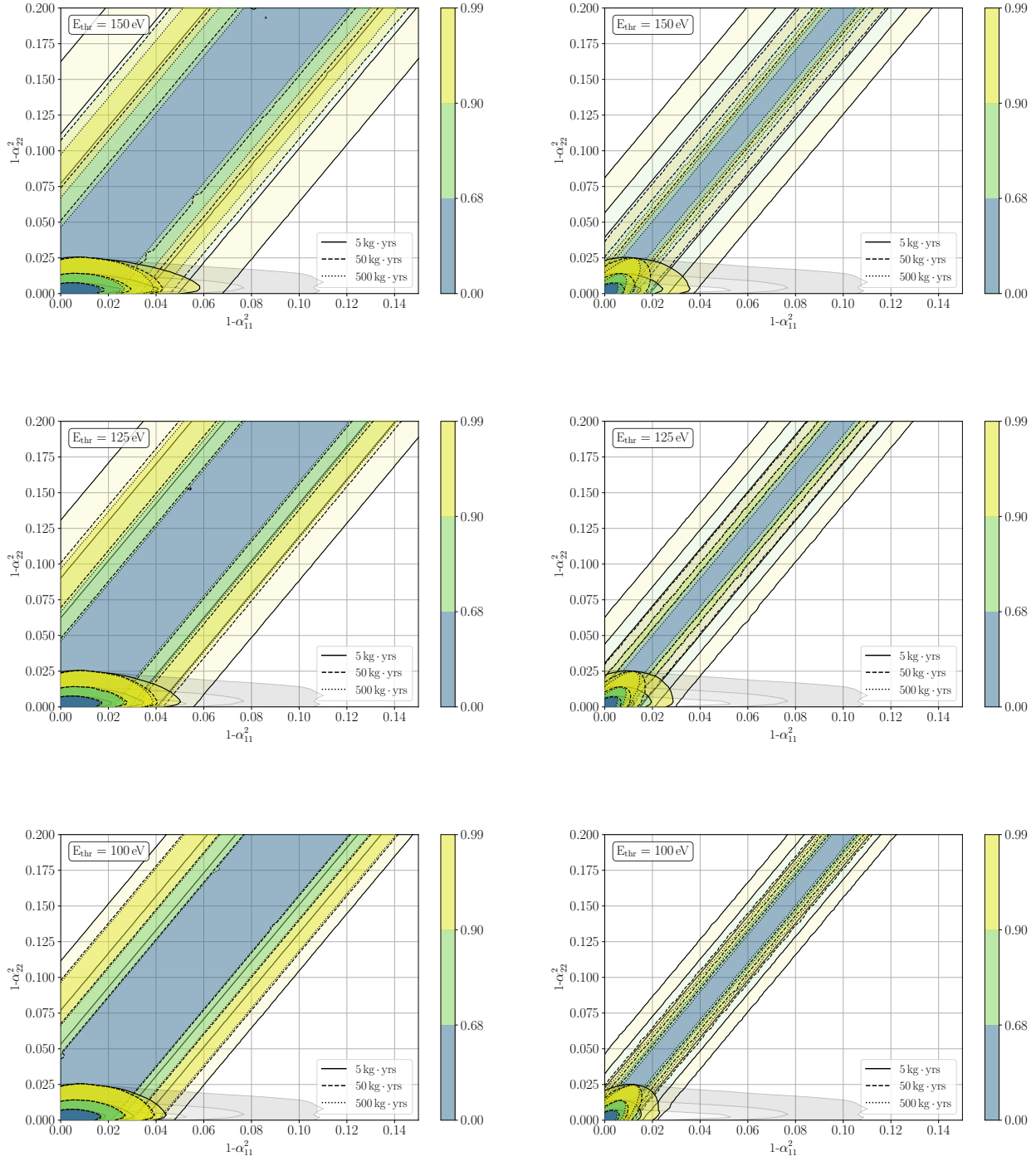


FIG. 7: Allowed regions of the alpha parameters for three threshold values and three exposures for our detector. The transparent contours show the limits of the assumed experimental configuration alone, while the non-transparent regions indicate parameter space still allowed when combined with knowledge from oscillation experiments. Gray contours show limits from oscillation experiments alone. Left: Results for our reference setup. Right: Results for improved knowledge on quenching (factor 2), background (factor 10) and reactor antineutrino flux (factor 10).

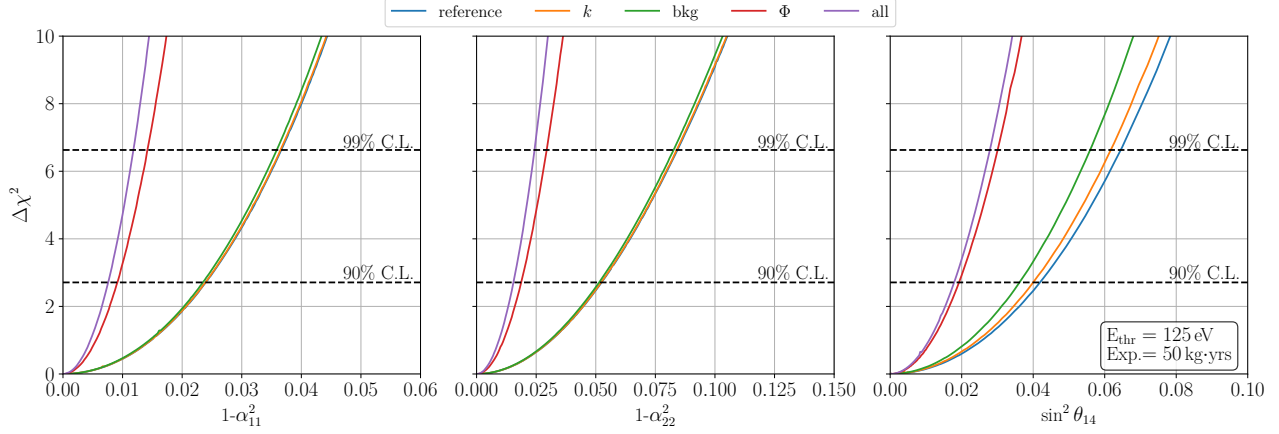


FIG. 8: $\Delta\chi^2$ profiles of individual model parameter for improved experimental characteristics: quenching uncertainty reduced by a factor 2 (orange), a factor 10 improvement on the background level (green) and flux uncertainty (red). The blue curve shows our reference configuration, while the purple lines indicate the combination of improvements. In the heavy seesaw limit (left and middle), the other alpha parameter is fixed to unity, while for the light sterile case (right) we set $\Delta m_{14}^2 = 1 \text{ eV}^2$ for illustrative purposes. As example a 125 eV-threshold detector with 50 kg·yr exposure is chosen.

The results of the full investigation are given in figure 7. As for the case of the single parameter investigation, we see the systematic uncertainties become limiting. For a threshold of 150 eV, there is a clear improvement between the exposures 5 kg·yrs and 50 kg·yrs. The soon available threshold value, i.e. 125 eV, already shows only a minor improvement between 50 kg·yrs and 500 kg·yrs, while for a 100 eV-threshold there is almost no increase in sensitivity. With the help of table I, it is possible to identify experimental configuration of almost similar sensitivity: 500 kg·yrs of exposure with a 150 eV-threshold is complementary to 50 kg·yrs exposure with a 100 eV-threshold. Such information are valuable and further development on the experimental site might decide which path to follow in the future.

The additional knowledge from oscillation sites - indicated by gray contours in figure 7 - strongly shrinks the parameter space still allowed by our CEνNS setup. As evident from the single parameter cases, limits on α_{22} are driven by oscillation experiments, while CEνNS experiments are contributing valuable knowledge for α_{11} , underlying their importance for global investigations.

The optimized configuration (a factor 10 improvement in flux uncertainty and background level as well as a factor 2 improvement on quenching parameter k) is able to constrain large parts of the parameter space. In particular, the transition from 5 kg·yrs to 50 kg·yrs shows a strong increase in sensitivity regardless of the chosen detection threshold. Note also that these systematics are still not limiting because improved thresholds and exposures lead to better limits on the non-unitarity parameters.

In order to quantify which experimental parameters are the main drivers of the improved sensitivity, we perform sensitivity estimates switching them on one after the other. Our findings are summarized in figures 8 for single parameters and 11 and 12 in the case of two parameters and both mass regimes. It is apparent that the uncertainty of the reactor antineutrino flux is the limiting factor, leading to a relative improvement of $\sim 63\%$ for the constraints on the non-unitarity parameters when reduced by a factor 10. Reducing the background level by a factor 10 only improves the limits by $\sim 2\%$. Further knowledge on the k parameter of the

Lindhard model seems to be of minor importance, i.e. $< 0.5\%$. Finally, an overall improvement of constraints on the alpha parameters of $\sim 70\%$ can be achieved when all factors are combined. Nevertheless, it is clear that better knowledge about reactor-related quantities will become the main limitation for the next-generation of experiment.

B. Light sterile limit

Similar to the previous case, we expect CE ν NS to have the dominating contribution to the achieved sensitivity, cf. figure 14 for exemplary single channel-sensitivities. In addition, it is worth noting that there already exist many dedicated experiments that aim to investigate a light sterile neutrino, especially at sites very close to a nuclear reactor core [91–95]. Our sensitivity results are illustrated in figure 9. We immediately see that in the light sterile limit systematic uncertainties are not limiting since higher exposures and lower detection thresholds still show constraining power. For example, a detector of 50 kg·yr exposure and a threshold of 125 eV mostly excludes mixing angles $\sin^2 2\theta_{14} \gtrsim 0.2$ for $\Delta m_{14}^2 \in [10^{-1}, 10] \text{ eV}^2$. The BEST experiment best-fit [96], $\Delta m^2 = 3.3_{-2.3}^{+\infty} \text{ eV}^2$ and $\sin^2 2\theta = 0.42_{-0.17}^{+0.15}$, is fully excluded in our projections, a result consistent with the strong tension between the BEST anomaly and other short-baseline data. Moreover, 100 eV-detectors with 500 kg·yr exposure will start to probe mixing angles $\sin^2 2\theta_{14} \lesssim 2 \cdot 10^{-2}$ in a setup at 20 m-distance. While in the previous case, configuration of almost similar sensitivity could be identified, i.e. 500 kg·yrs at 150 eV vs. 50 kg·yrs at 100 eV, the message here is different: CE ν NS searches for a light sterile neutrino clearly benefit from a lower detection threshold.

Improved experimental specification, i.e. antineutrino flux, background level and quenching, also boost the experimental sensitivity in this context. In particular, a setup with a threshold below 125 eV and an exposure larger than 50 kg·yrs will clearly exclude mixing angles above 0.1, and larger parts above 0.05.

To quantify the effect of experimental uncertainties also for this case, we fixed $\Delta m^2 = 1 \text{ eV}^2$ and varied them individually.⁴ The $\Delta\chi^2$ profile in terms of the mixing angle $\sin^2 \theta_{14}$ are given in figure 8 (right plot). Also here, the flux uncertainty is the driving factor of the obtained sensitivity with $\sim 54\%$ relative improvement. The impact of quenching and background level is stronger, i.e. $\sim 4\%$ and $\sim 14\%$ relative improvement, but are still of secondary importance. An overall improvement of $\sim 57\%$ can be gained with the combination of these factors, less strong than for the seesaw limit. The effects in the two-dimensional parameter space are illustrated in figure 12.

However, the parameter space probed is mostly excluded by existing short-baseline experiments. The advantage of CE ν NS setups in this context lies in their compactness. An experiment of several medium-size CE ν NS detectors at different distances would allow to reduce reactor-related uncertainties. In addition, CE ν NS is sensitive to all neutrino flavors, providing complementary information to charged current electron-flavor-based searches.

⁴ Here we have selected a generic point close to the second main oscillation peak. Of course, the results vary depending on the chosen value of Δm^2 and the chosen experimental characteristics.

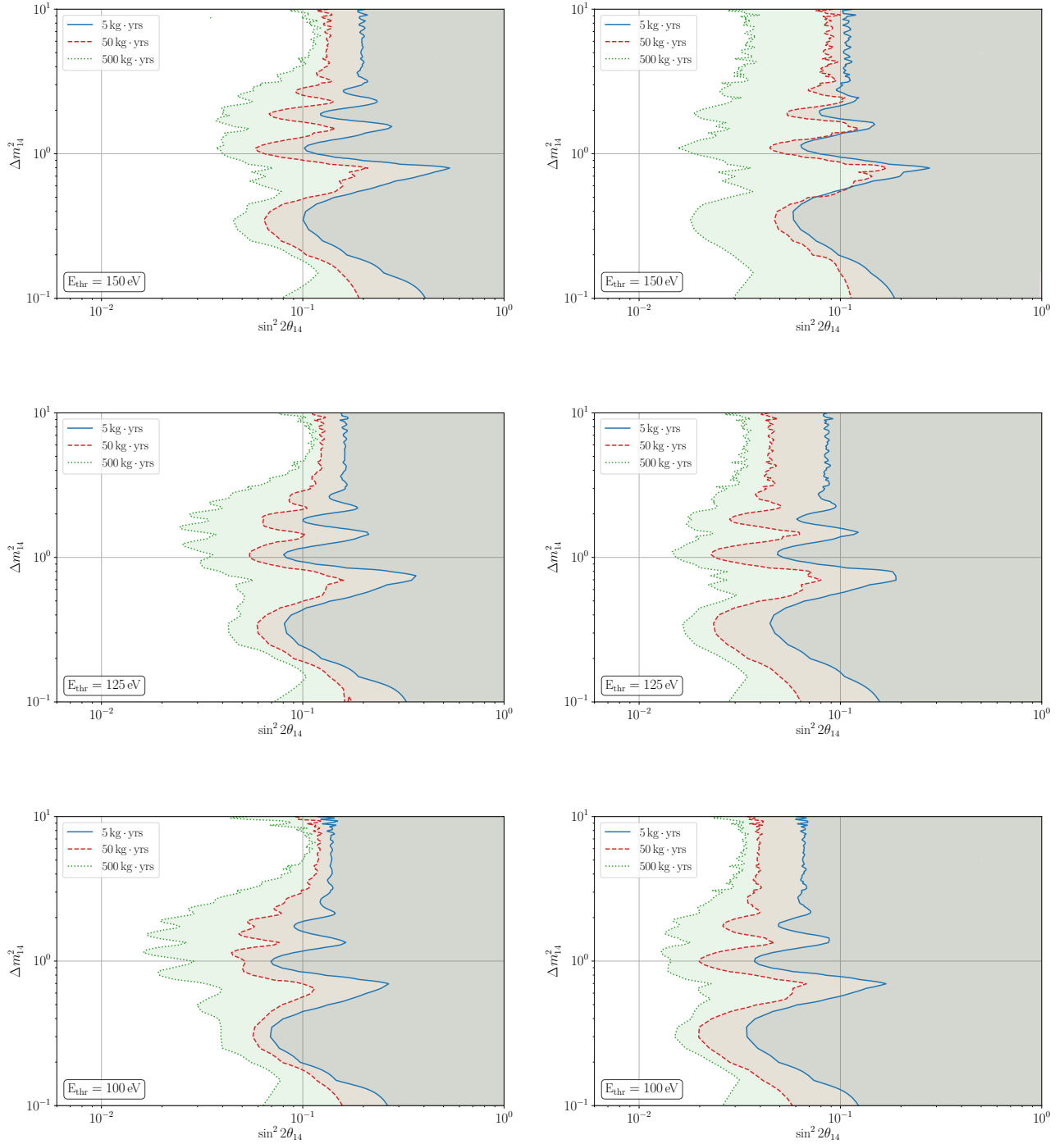


FIG. 9: Experimental sensitivity (exclusion potential) of light sterile neutrino searches of our reference setup (left column) with the assumed detection thresholds (150, 100, 50) eV and exposures (5, 50, 500) kg.yrs. In contrast to the previous case, such searches are not limited by the systematic uncertainties assumed. Results for improved experimental characteristics (a factor 10 improvement in flux uncertainty and background level as well as a factor 2 improvement in quenching uncertainty) are shown as well (right columns).

	M limit from α_{11}	M limit from α_{22}
oscillations	700 GeV	1740 GeV
oscillations + now	1300 GeV	1940 GeV
oscillations + soon	1330 GeV	1940 GeV
oscillations + future	2460 GeV	2250 GeV

TABLE III: Limits (at 90% C.L.) on the heavy mediator mass scale M inferred from the non-unitarity parameters α_{11} and α_{22} , assuming a low-scale seesaw with $\mathcal{O}(1)$ coefficients. The row “oscillations” uses current bounds from Ref. [80]. The rows “oscillations + current/realistic/optimistic” additionally include projected CE ν NS sensitivities for CONUS-like detector thresholds and exposures of 150 eV with 50 kg yr (now), 125 eV with 50 kg yr (soon), and 100 eV with 500 kg yr (future). In the optimistic case we assume the optimized experimental configuration, see the main text for more details.

V. CONCLUSIONS

After the first detection of CE ν NS, more and more CE ν NS experiments start data collection and we can expect further interesting physics results in the future, not only from experiments but also from subsequent phenomenological studies. The existing data and analyses already revealed the large potential of CE ν NS that could be further exploited when the next generation of experiments transitions to precision measurements. Motivated by the recent CE ν NS observation of the CONUS+ experiment, the present work aimed at identifying the sensitivity of a future upgrade based on demonstrated technology, i.e. a next-generation Germanium-based experiment close to a nuclear power reactor site. With lepton non-unitarity as example, the expected experimental reach has been determined for well-selected benchmark points, while taking into account major experimental uncertainties such as reactor antineutrino flux and signal quenching. Further, the impact of these uncertainties and the underlying background level in future experimental endeavors has been assessed to identify key drivers for scientific progress in this context.

In the so-called seesaw limit, where new degrees of freedom are heavy, a future CONUS-like experiments contribute valuable information on the non-unitarity parameters, cf. table I and figures 5 and 7. A CE ν NS setup with characteristics soon to be achieved (50 kg·yr, 125 eV-threshold) will be able to probe non-unitarity parameters related to energy scales of 1100 GeV (for α_{11}) and 760 GeV (for α_{22}). With reduced uncertainties on the reactor antineutrino flux (factor 10), quenching (factor 2) and a reduced background level (factor 10) scales up to 2500 GeV (for α_{11}) and 1700 GeV (for α_{22}) could be tested in a future setup. In general, when combined with knowledge from oscillation experiments even higher scales can be probed, see table III. We emphasize that the mass scales inferred from Eq. 9 are not strict experimental limits: they rely on the assumptions specified in the main text, including the low-scale seesaw interpretation and $\mathcal{O}(1)$ coefficients.

For the case of light new particle (light sterile limit), bounds from CE ν NS experiment will be improved significantly, cf. figure 9. Especially, mixing angles above ~ 0.1 could be fully excluded when systematic uncertainties and background can be further lowered. Obtained results are not competitive to existing bounds, but are exceptional in the sense that they are flavor-independent. Further, more refined investigation with several CE ν NS detectors may be envisioned. For future design, our work contributes interesting knowledge since experimental configurations with (roughly) the same sensitivity have been identified. In particular, the question whether to build an experiment with 500 kg·yr exposure and a 150 eV-threshold or 50 kg·yr with a 100 eV-threshold may be answered by detector developments in the next years. Furthermore, it has become clear that

systematic uncertainties will be the limiting factor in future precision experiments underlining the importance of carefully assessing an experiment's uncertainties and improved theory predictions. In the context of this work, the reactor antineutrino flux is identified as one of the key drivers for experimental sensitivity, cf. figure 8. Of course, more refined studies from our experimental colleagues are needed to fully consider all potential uncertainties underlying a specific experimental setup.

Nevertheless, our work clearly underlines the strong potential of future CE ν NS experiments for future tests of the lepton sector and searches of physics beyond the standard model.

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Appendix A: Detailed $\Delta\chi^2$ profiles of the seesaw limit

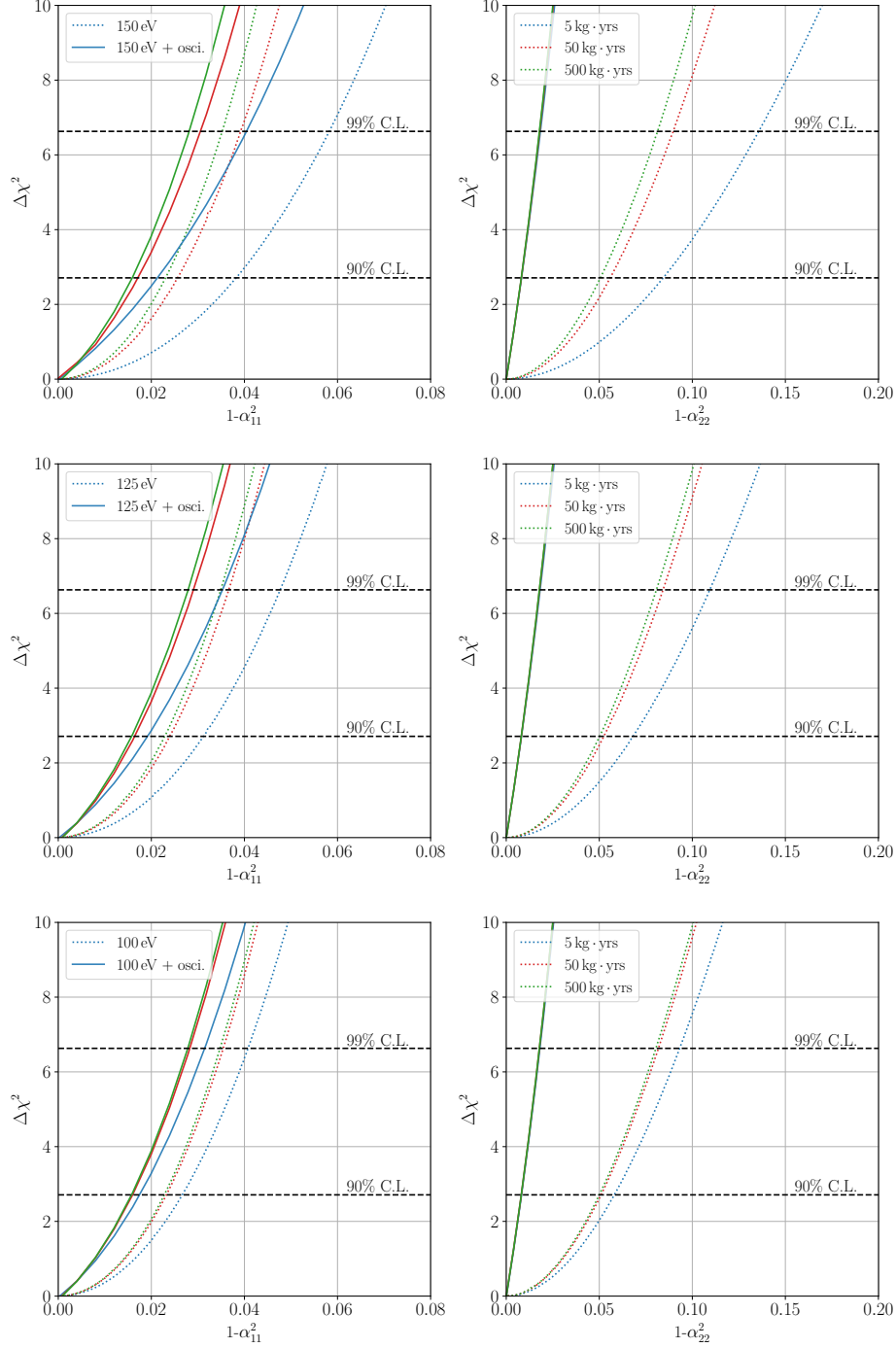


FIG. 10: Detailed $\Delta\chi^2$ profiles for the individual alpha parameters and the three threshold values (150 eV, 125 eV, 100 eV from top to bottom) under consideration. The experimental reach of the individual experimental configuration (dashed lines) is shown in comparison to the sensitivity when information from oscillation experiments is added as external knowledge to our analysis (solid).

Appendix B: Impact of improved experimental parameters - full parameter space

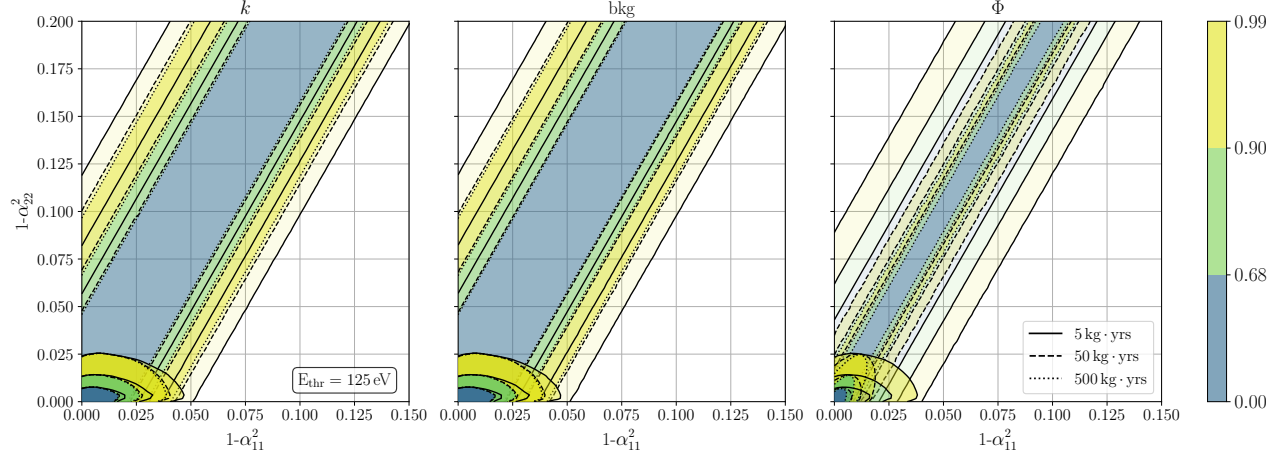


FIG. 11: Limits on the non-unitarity parameters under optimization of experimental characteristics: quenching uncertainty reduced by a factor 2 (left), background reduction by a factor of 10 (middle) and antineutrino flux uncertainty reduced by a factor 10 (right). Here, a 125eV-threshold detector with 50 kg·yr exposure is chosen as example.

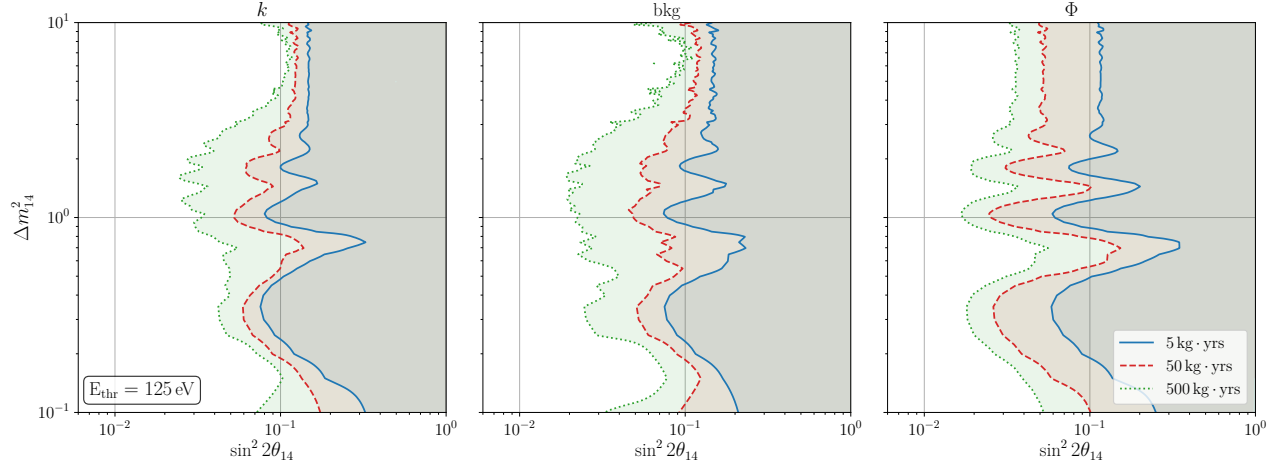


FIG. 12: Limits on mass and mixing parameters under optimization of experimental characteristics: quenching uncertainty reduced by a factor 2 (left), background reduction by a factor of 10 (middle) and antineutrino flux uncertainty reduced by a factor 10 (right). Here, a 125 eV-threshold detector with 50 kg·yr exposure is chosen as example.

Appendix C: Single channel comparison

1. Seesaw limit

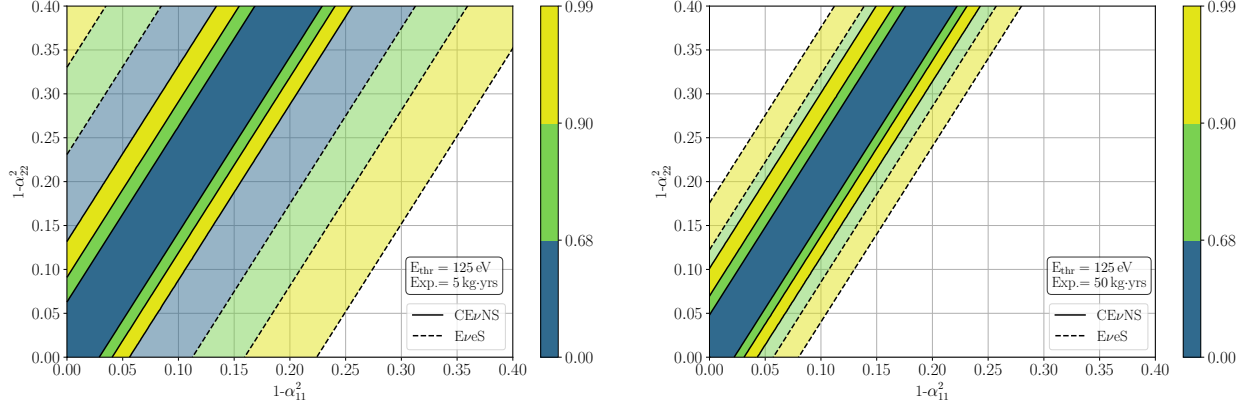


FIG. 13: Sensitivity of $\text{CE}\nu\text{NS}$ and $\text{E}\nu\text{eS}$ for the assumed experimental configuration. Limits (at 90% C.L.) on the alpha parameters are given for $\text{CE}\nu\text{NS}$ below 1 keV and $\text{E}\nu\text{eS}$ above 1 keV up to 100 keV. It is evident that the obtained results are dominated by $\text{CE}\nu\text{NS}$.

2. Light sterile limit

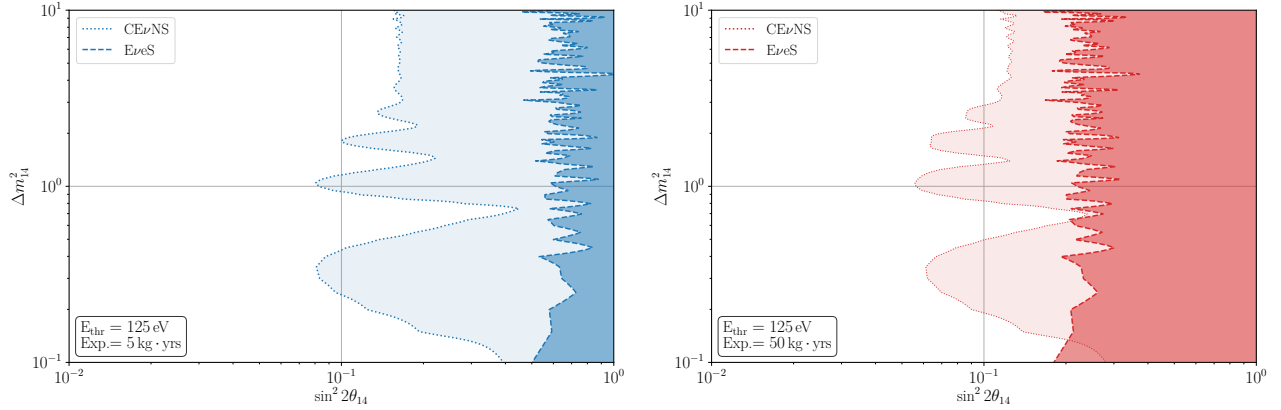


FIG. 14: Exclusion potential of $\text{CE}\nu\text{NS}$ and $\text{E}\nu\text{eS}$ for the assumed experimental configuration. Limits (at 90% C.L.) on the mixing angle and the mass-squared difference are given for $\text{CE}\nu\text{NS}$ below 1 keV and $\text{E}\nu\text{eS}$ above 1 keV up to 100 keV. Also here, $\text{CE}\nu\text{NS}$ is the determining factor for the expected sensitivity.